A major purpose of the Technical Information Center is to provide the broadest dissemination possible of information contained in DOE's Research and Development Reports to business, industry, the academic community, and federal, state and local governments.

Although a small portion of this report is not reproducible, it is being made available to expedite the availability of information on the research discussed herein.

1

LA-UR-84-1873

Los Alamos National Laboratory is operated by the University of California for the United States Department of Energy under contract W-7405-ENG-36

LA-UR--84-1873

DE84 013903

TITLE: QUARK TUNNELING IN HYPERNUCLEI

AUTHOR(S): T. Goldman, T-5

SUBMITTED TO For proceedings of the Conference on the Intersections Between Particle and Nuclear Physics, Steamboat Springs, CO, May 23-30, 1984

MASTER

Los Alamos National Laboratory
Los Alamos, New Mexico 87545

FORM NO 836 R4
ST NO 2629 5/81

QUARK TUNNELING IN HYPERNUCLEI

T. Goldman
Theoretical Divison, Los Alamos National Laboratory
Los Alamos, New Mexico 87545

ABSTRACT

A new model of nuclei is described. Its qualitative successes are discussed. Its implications for and new tests in hypernuclei are introduced.

I begin from the assumption, based on its successes in particle physics and hadron spectroscopy,[1] that Quantum Chromo-Dynamics (QCD) is the correct theory of the strong interactions. Nuclear physics must then be a derivable aspect of QCD and one must ask how to achieve this. Conventional nuclear physics implicitly assumes that one should first construct baryons and mesons, and then analyze their residual, color singlet (van der Waals-like) interactions. Although this has been a very successful phenomenology, it has no mathematical justification within QCD.

To see this, consider the QCD path integral with quark sources $\eta,\bar{\eta}$:

$$Z[\eta,\bar{\eta}] = \int d[q]d[\bar{q}]d[G] \exp i \{S_{QCD} + \int d^4x(\bar{q}\eta + \bar{\eta}q)\} \tag{1}$$

where S_{QCD} is the QCD action and G represents gluon fields, and the functional derivative

$$\left.\frac{\delta^{6A}Z}{\delta\bar{\eta}_1 \ldots \delta\bar{\eta}_{3A}\delta\eta_1 \ldots \delta\eta_{3A}}\right|_{\eta,\bar{\eta}=0} = \langle 0|\bar{q}_1 \ldots \bar{q}_{3A} q_1 \ldots q_{3A} e^{iS_{QCD}}|0\rangle \tag{2}$$

This describes, at once,

a) free propagation of A nucleons

b) 2, 3, . . . A-1 body interactions of A nucleons with otherwise free propagation ·

c) A-nucleon scattering resonance and bound state poles.

Part c) includes all nuclei of baryon number A and their excited states. (Part b) includes all smaller nuclei.) In Eq. (2), all gluon fluctuations have been summed over, and all higher quark Fock components are implicit. In principle, even strange (s-$\bar{s}$), charm (c-$\bar{c}$), etc. nuclear components are included.

At present, we do not know how to evaluate Eq. (2) directly and non-perturbatively. Even lattice QCD calculations,[2] which study hadron spectroscopy, are presently limited to lattices with an overall size of about 1 fm. So we must make a guess as to what classical configurations dominate the path integral. The guess which is the starting point of conventional nuclear physics is that the 3A quarks clump into A regions of relatively high quark density, which we identify as nucleons, and that there are no significant phase relations between the quark amplitudes in these regions.

However, the motion of the quarks relative to the c. m. of each nucleon implies that the quark wavefunctions, even from nucleons in different orbitals, are not orthogonal. Further, the meson exchange picture of the nuclear force populates the inter-nucleon regions with finite quark amplitudes. Finally, the nucleon size is large compared to the average inter-nucleon spacing, whether evaluated for a full Breuckner-Hartee-Fock (BHF) A-body wavefunction, or for close packed hard spheres. All of these points lead us to believe that phase correlations can, and will, develop between quark waves nominally in different nucleons.

I want to emphasize the unreasonableness of omitting such phase relations, *ab initio*. Although the quark amplitude may rapidly diminish in the region beyond 1 fm from the c. m. of a hadron, we know on general quantum mechanical grounds that the overall quark eigenstates of two such hadrons, <u>at any separation,</u> must be the symmetric and antisymmetric combinations of the individual wavefunctions. For "large" separations, these two states will be almost degenerate, and we can ignore the difference for "short" interaction times. But "large" and "short" should be demonstrated. It could happen that the effects of the phase correlations are negligible. Our calculations, however, show that, at nuclear matter density, the internucleon separations are <u>not</u> large enough, and the nuclear time scale is <u>not</u> short enough to ignore these phase correlations. They have significant effects on the scale of nuclear binding energies. This is a common effect in solid state physics and in chemistry. Two examples: The electrons in the bands contribute to the stiffening of crystals. The benzene ring is not made of Carbon atoms, but of Carbon "cores"; the binding and stiffening are due to delocalized, phase coherent electrons.

To proceed, we must give an ansatz, ultimately to be tested from Eqs. (1) and (2), which reflects the lumpiness of the quark density distribution in nuclei. We assume[3] that the result of the path integration over the gluon field and over the quark-antiquark pair fluctuations is an array of A truncated QCD potentials, the eigenstate of which are to be filled by 3A quarks. The QCD potential is that between a quark (color triplet) and an anti-triplet of color. The A origins for this confining potential, which we term color singlet centers (CSC), correspond to nucleon c. m.'s for an intrinsic BHF wavefunction. The potential at a given point is that calculated from the nearest CSC; thus the potentials are truncated in the nuclear interior. Quarks are still confined to the region of the nucleus by the potential arising from the CSC's on the nuclear surface. Fig. 1 shows a one dimensional representation of our mean field potential and the lowest eigenmode.

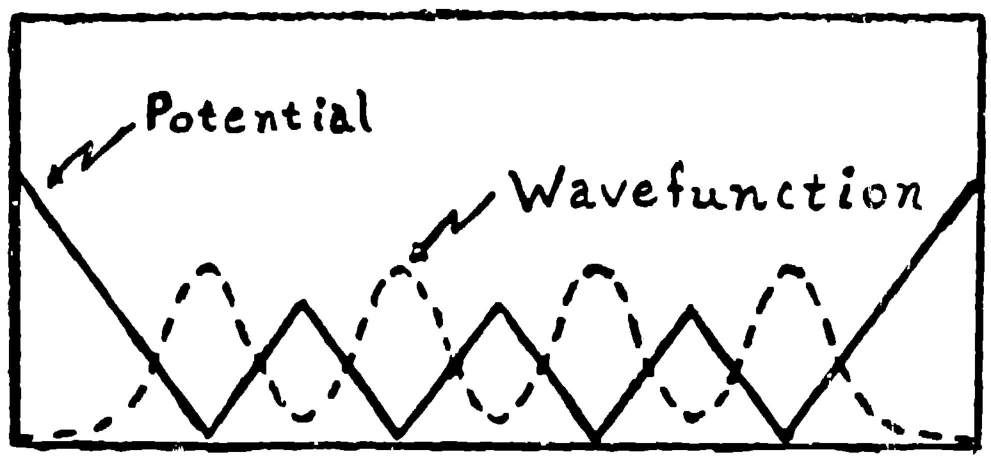

Fig. 1. Mean field potential and wavefunction of lowest eigenmode.

Before embarking on detailed calculations of nuclei in this model, we examined a toy problem[3] to see if anything like real nuclei resulted from our ansatz. This toy problem involved cubical nucleons, with delta-functon barriers on the faces between cubes, and an infinite barrier at the surface. The delta-function barrier strength increased as the square of the separation between the nucleons. The advantages of the Cartesian separation of variables are obvious.

The toy problem exhibits the expected features: The lowest eigenstate has a quark energy lower than in an individual (cubical) nucleon. The highest filled state has the same wave number as in a nucleon, but oscillates coherently across the nucleus (180° phase change from one cube to the next). The wavefunctions for A=27 are shown in Fig. 2. The overall energy of the nucleus rises as it is compressed, due to stronger localization of the quark waves. The overall energy also rises if one increases the separation between the nucleons due the increased barrier strength and so, reduced tunneling. The mean binding energy per nucleon saturates at 30 - 60 MeV, for large nuclei and reasonable values of the barrier strength.

Using West's technique,[4] we have also found the quark probability distribution in our nuclei. These show enhancements both at low and high momentum fraction, x, relative to the distribution in a nucleon. The former, known as the EMC effect, is due to quark delocalization in our nuclei, and the latter is due to collective nuclear recoil, often described as Fermi motion. We find a weak increase in the EMC effect[5] with increasing A, qualitatively consistent with the data. We have not attempted a detailed fit to the data using this highly artificial toy model.

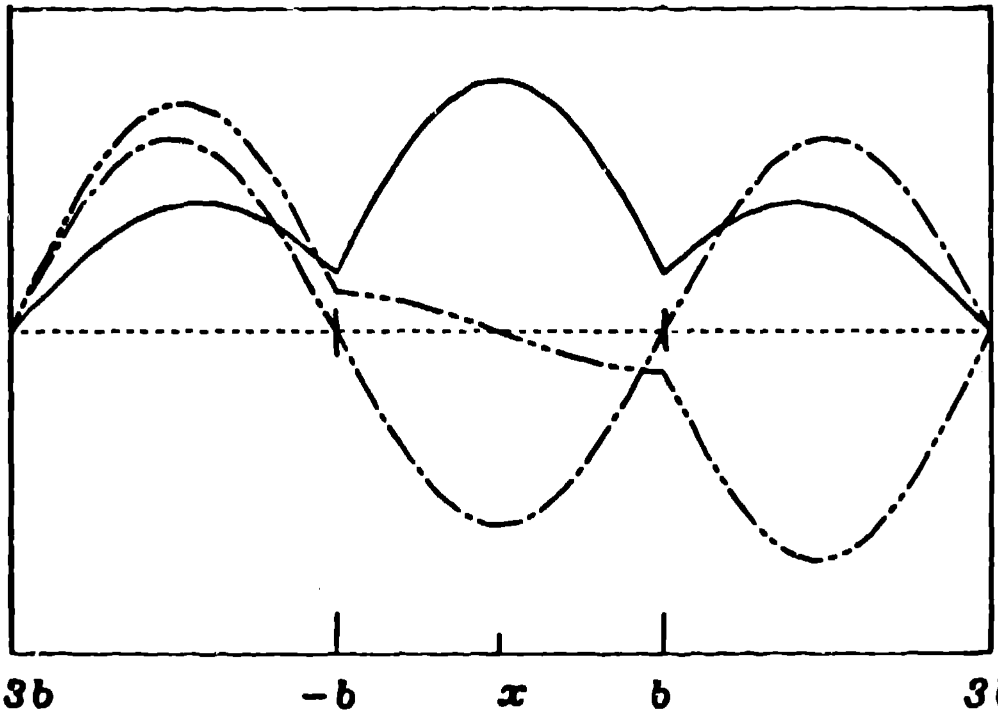

Fig. 2. Wavefunctions of three lowest Cartesian eigenmodes for A=27; b=0.93 fm.

The toy problem illuminates many features of the true model. One is led naturally to a quark "shell" model of nuclei. The energy eigenstates cluster, forming almost a band, ranging from 60 to 70 MeV below the eigenenergy of a quark in a nucleon. A second "band", with wider spacing of the modes, appears from about 100 MeV above the top to twice the top energy of the lowest band. It is clear from this, and Pauli exclusion, that the quark higher Fock components are suppressed: The addition of an antiquark provides much less binding energy than must be paid to be put a quark in the second band. Of course, the intrinsic wavefunctions must be projected to definite angular momentum states, which we did not do in the toy problem.

It is immediate in this quark "shell" model that double charge exchange (DCX) in πA scattering is not a convolution of two single charge exchanges (CEX). CEX proceeds dominantly by u $\leftrightarrow$ d quark interchange, producing a strong η signal[6] when phase space is not a constraint, while DCX requires $\bar{d}d$ annihilation into $\bar{u}u$ to accompany CEX. Thus, two u's must contend with Pauli blocking. Strangeness exchange is similarly straightforward: The s-quark in a K^- may exchange with a u from the target nucleus, but the $\bar{s}$ from a K^+, to first order, finds no appropriate target, and so is suppressed.

In the quark "shell" model, hypernuclei appear to be very different than in the conventional picture. Instead of a Λ sinking down to the lowest nucleon state, it is ripped apart: The strange quark may settle into the bottom of the well, but its accompanying up and down quarks are blocked, and must occupy the highest levels in the nucleus. If viewed as a single object, the Λ stays at a higher average level than expected. The net result is a smaller binding energy for the Λ than expected.[7]

Hypernuclei provide a wonderful new testing ground for our quark model of nuclei. Unlike conventional calculations, which acquire myriads of essentially free parameters which must be measured, (although in principle, they can be calculated from QCD) such as the Λ-N scattering length and Λ spin-orbit coupling, our model has only the strange quark mass as a new parameter, and this must be consistent with the Λ, Ω, etc. masses. Having fixed this, we can predict hypernuclear masses with no free parameters.

Although we have not yet done so with precise calculations, simple, parton-model-like estimates can be made quickly. In ^{4}He, we must replace one of the 12 ordinary quarks with a strange one. Because it is massive, delocalization is less significant for it and so it provides less binding energy: We estimate

$$^4_\Lambda\mathrm{He} \cong {}^4\mathrm{He} + \sqrt{E_q^2 + m_s^2} - E_q \tag{3}$$

The binding energy is about 20 MeV.

The result is quite sensitive to m_s, the strange quark mass. A more precise quark-nucleus calculation should yield a significant value of m_s from the experimental value of the binding energy. As it is, the same parameters as in Eq. (3) imply that the binding energy of $^5_\Lambda$He is 45 MeV. These estimates do not include any reduction due to the separation of the s-quark from its u,d partners in different states. This is also true for the corresponding 13 MeV estimated binding energy in large nuclei.

Although these results are not accurate, we have landed in the right ballpark with an extremely crude approximation. Conventional calculations must invoke Λ-Σ coupling (mixing) to avoid overbinding[7] the Λ. In our picture, $^A_\Lambda Z$ is actually a misnomer: An ordinary quark has been replaced by a strange quark, but there is nothing to indicate any organization of quarks into a Λ or a Σ. Only the overall isospin is well defined; as such, Λ-Σ mixing is to be expected, as both can contribute to a state of given isospin and strangeness.

As an aside, we raise the question of "hyponuclei," namely $^A Z + \bar{\Lambda}$. Here the ordinary antiquarks as well as the (anti) strange quark can sink to the lowest energy state. Also, there is no additional color electric energy required to separate the three antiquarks into different eigenstates as there was for the quarks. Thus, we expect significantly larger binding. Unfortunately, it will not be large enough to prevent

$$ {}^{A-2}_{\bar{\Lambda}}Z \rightarrow {}^{A-2}Z + K + n\pi\text{'s} \qquad (4) $$

which is an open channel with at least 1 GeV available. Observing the binding energy of such a short-lived hyponucleus is probably very difficult.

The last test we would like to discuss is hypernuclear beta decay. As you well know, $\Lambda \rightarrow N + \pi$ in a large nucleus is Pauli blocked, but the weak scattering $\Lambda N \rightarrow NN$ may still occur. Our model draws the same conclusions on the quark level: $s \rightarrow ud\bar{u}$ is blocked but $su \rightarrow ud$ can occur. In ${}^4_\Lambda He \rightarrow {}^4He + \pi$, the channel is even open and the wavefunction overlap is good. In ${}^5_\Lambda He$, the decay channel with good overlap is blocked, but there is sufficient phase space to reach "${}^5He$", where the overlap is poor.

Semileptonic decays may provide more distinctive and quantitative tests. Since the weak decay varies as the fifth power of the energy released, we expect at least a 10% difference from the absolute semi-leptonic decay rate of a free Λ. It should be possible to make this estimate very accurate, even beyond the spectator approximation. The result should be different from conventional calculations because the strange quark provides virtually all of the binding energy. In the conventional picture, this is shared with the u and d quarks since it is the Λ, as a whole, which binds.

In conclusion, I have described a potentially viable new model for nuclear physics, which is closely linked to QCD and has essentially no free parameters. Although I have not described a realistic application, I am presently working with K. Schmidt and G. J. Stephenson, Jr. on ${}^4He$, using the full Dirac equation with QCD potentials derived from charmonium and upsilon studies. So far, we have found that quark delocalization provides at least as much binding energy as needed to match experiment.

I have indicated that spectra and decay rates in hypernuclei may be predicted to be very different from what is expected in the conventional picture. This is in addition to new observations applicable to ordinary nuclei, such as the implicitly phonon character of excitation spectra, and the virtual absence of pionic field in any nucleus.[8] Because of the significant differences between a strange quark and a point Λ in a nucleus, hypernuclear physics offers an outstanding opportunity to clarify one of the forefront problems in nuclear physics today - namely its relation to QCD.

Let me close with the observation that if this new picture of nuclei is experimentally confirmed, the vast wealth of nuclear phenomena will then foster the study of QCD at large distances.

FOOTNOTES AND REFERENCES

1. See, for example, C. Quigg and J. L. Rosner, Phys. Rept. 56, 167 (1979).
2. The need for small lattice link sizes has been shown by G. A. Baker, L. P. Benhofy, F. Cooper and D. Preston, Nucl. Phys. B 210, 273 (1982). Our statement follows from the practical limitation of ~10 links per dimension in 3- and 4-dimensional systems.
3. T. Goldman and G. J. Stephenson, Jr., "Quark Tunneling in Nuclei," Los Alamos preprint LA-UR-84-1645.
4. G. B. West, Ann. Phys. (N. Y.) 74, 464 (1972); W. B. Atwood and G. B. West, Phys. Rev. D 7, 773 (1973).
5. J. J. Aubert et al., Phys. Lett. 123B, 275 (1983); A. Bodek et al., Phys. Rev. Lett. 51, 534 (1983).
6. Experiment 852 (J. C. Peng, spokesman) has been approved by the LAMPF PAC to study η production in pion scattering.
7. See, for example, the contribution to this parallel session by B. F. Gibson.
8. There is already some experimental support for this position; see T. A. Carey et al., Los Alamos preprint LA-UR-84-1207, submitted to Phys. Rev. Lett.